

# Exotic baryons from a heavy meson and a nucleon

## – Negative parity states –

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### Abstract

We study heavy baryons with an exotic flavor quantum number formed by a heavy meson and a nucleon ( $\bar{D}N$  and  $BN$ ) through a long range one pion exchange interaction. The bound state found previously in the  $(I, J^P) = (0, 1/2^-)$  channel survives when short range interaction is included. In addition, we find a resonant state with  $(I, J^P) = (0, 3/2^-)$  as a Feshbach resonance predominated by a heavy vector meson and a nucleon ( $\bar{D}^*N$  and  $B^*N$ ). We find that these exotic states exist for the charm and heavier flavor region.

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## I. INTRODUCTION

Hadron physics has opened a renewed interest in multi-hadron systems. The most familiar example is the atomic nucleus which is the bound state of protons and neutrons. However, interestingly enough, we do not have yet clear evidences for analogous systems of baryon number one or zero. Yet a well-known candidate is  $\Lambda(1405)$  which is considered to be a quasi-bound state of  $\bar{K}N$  and  $\pi\Sigma$  [1–5]. The isoscalar meson which has been established as a resonance in pseudoscalar meson scatterings may also be a quasi-bound state of the two mesons [6–9]. Such hadronic composites, or molecular states, are dynamically generated via hadron-hadron interactions, and are expected to appear in various mesonic and baryonic systems. The hadron composites are important also for study of the hadron dynamics in nuclear matter [10–15].

In the constituent quark model, the generation of hadronic composites can be understood as formation of clusters in multi-quark systems. Because a typical excitation energy of hadron resonances amounts to several hundred MeV and is enough to create a  $\bar{q}q$  pair of constituent quarks, a multi-quark component naturally appears in the resonance states in addition to the minimal configuration of  $\bar{q}q$  or  $qqq$  with an orbital excitation. Such a multi-quark configuration may arrange itself into a set of color-singlet clusters, namely a set of hadrons. This serves a quark model picture of hadronic composites.

Recently, a novel structure has been suggested by one of the present authors in manifestly exotic channels with one (anti) heavy quark, for example,  $\bar{D}N$  whose minimum quark content is  $uudd\bar{c}$  [16]. This is a charm analogue of the pentaquark  $\Theta^+ \sim uudd\bar{s}$  [17, 18]. So far the charmed pentaquarks in various forms have been discussed by many authors [19–29]. To obtain a bound  $\bar{D}N$  system composed of a  $\bar{D}$  meson and a nucleon, the one pion exchange interaction was found to be crucially important. As emphasized in Ref. [16], in particular, the tensor force yields strong attraction through the mixing of an  $S$ -wave state of  $\bar{D}N$  and a  $D$ -wave state of  $\bar{D}^*N$ . The mixing effect is more important for heavier flavor sectors, where pseudoscalar and vector mesons are more degenerate. Indeed the mass splittings of  $\bar{D}\bar{D}^*$  and of  $BB^*$  are about 140 and 46 MeV, respectively, and are much smaller than the one of the strangeness sector ( $KK^*$  mass difference  $\sim 400$  MeV).

Such a mixing mechanism of  $S$ - and  $D$ -waves (more generally mixing of  $L$ - and  $(L \pm 2)$ -waves, with  $L$  being an orbital angular momentum) has been known to be very important

for the deuteron binding, while its relevance has been reexamined for other nuclear systems rather recently [30, 31]. In QCD, spontaneous breaking of chiral symmetry is underlying; the light pseudoscalar mesons, pions, are generated as the Nambu-Goldstone modes, with strong coupling to hadrons [32, 33]. It is emphasized that the pseudoscalar nature of the pion necessarily leads to the Yukawa coupling of the form  $\vec{s} \cdot \vec{q}$ , where  $\vec{s}$  is the spin operator for the particles coupling to the pion and  $\vec{q}$  is the momentum of the pion. Then this coupling provides the tensor force in two-body systems. It is noticeable that such one pion exchange interaction can lead also to stable exotic states of two heavy meson systems as pointed out in Refs. [26, 34–36]. In the present study, we investigate the exotic state of a heavy meson and a nucleon, namely  $\bar{D}N$  and  $BN$ .

In this paper, we examine the system of  $PN$  and  $P^*N$  with the inclusion of short range interactions. Here and in what follows we introduce the notation  $P(= \bar{D}, B)$  for a heavy pseudoscalar meson and  $P^*(= \bar{D}^*, B^*)$  for a heavy vector meson. We study not only bound states but also resonances. We concentrate our analysis on the low lying states in which  $S$ -wave component is included. Furthermore, we present our analysis only in isospin singlet channels because we find neither bound state nor resonant state in isospin triplet channels.

In section II, we briefly describe the interactions between  $PN$  and  $P^*N$  based on the heavy quark symmetry. It has been known that the heavy quark symmetry plays an important role for charm and bottom quarks, not only in the dynamics of quarks [37–39], but also in the dynamics of heavy hadrons [39–44]. Therefore, in the present study, we employ the interaction from heavy quark symmetry. For comparison, we also investigate the interaction derived from flavor  $SU(4)$  symmetry [45–48]. In this paper, we include not only the pion exchange but also vector meson exchange interactions to show the dominant role of the pion exchange interaction. We discuss bound states and scattering states in sections III and IV, respectively. We confirm that a bound state exists for isospin, spin and parity  $(I, J^P) = (0, 1/2^-)$ . For resonance, we find a new state in the  $(I, J^P) = (0, 3/2^-)$  channel having a narrow decay width as a Feshbach resonance predominated by  $P^*N$ . In section V, we discuss flavor dependence of the present results by varying the mass of the heavy mesons continuously, and show that the existence of the above exotic states is the feature of heavy flavors. Final section is devoted to summary and discussions.

TABLE I: Various coupled channels for a given quantum number  $J^P$  for negative parity  $P = -1$ .

$J^P$	channels
$1/2^-$	$PN(^2S_{1/2})$ $P^*N(^2S_{1/2})$ $P^*N(^4D_{1/2})$
$3/2^-$	$PN(^2D_{3/2})$ $P^*N(^4S_{3/2})$ $P^*N(^4D_{3/2})$ $P^*N(^2D_{3/2})$

## II. INTERACTIONS

The systems we are interested in have a manifestly exotic flavor structure of  $qqqq\bar{Q}$  in its minimal quark content, which is a heavy quark analogue of the pentaquark  $\Theta^+ \sim qqqq\bar{s}$ . Here  $Q$  and  $q$  denote heavy and light ( $u, d$ ) quarks. We investigate whether these exotic baryons are formed as (quasi) bound states of a  $\bar{Q}q$  meson (denoted by  $P$  or  $P^*$  following introduction) and a  $qqq$  nucleon (denoted by  $N$ ). For this picture to work well, these two hadronic ingredients should be sufficiently apart and keep their identities. Possible effects of internal structure is then expressed by form factors. The two-body states of a pseudoscalar (vector) meson and a nucleon are classified by their isospin  $I$ , total spin  $J$  and parity  $P$ , or orbital angular momentum  $L$ , where  $P = (-1)^{L+1}$ . For a given  $J^P$ , there are three or four coupled channels as summarized in Table. I, where low lying components including  $S$ -states are shown.

To obtain interactions for heavy mesons and nucleons, we employ Lagrangians satisfying heavy quark symmetry and chiral symmetry [49]. They are well-known and given as

$$\mathcal{L}_{\pi HH} = ig_\pi \text{Tr} [H_b \gamma_\mu \gamma_5 A_{ba}^\mu \bar{H}_a] , \quad (1)$$

$$\mathcal{L}_{v HH} = -i\beta \text{Tr} [H_b v^\mu (\rho_\mu)_{ba} \bar{H}_a] + i\lambda \text{Tr} [H_b \sigma^{\mu\nu} F_{\mu\nu}(\rho)_{ba} \bar{H}_a] , \quad (2)$$

where the subscripts  $\pi$  and  $v$  are for the pion and vector meson ( $\rho$  and  $\omega$ ) interactions, and  $v^\mu$  is the four-velocity of a heavy quark. In Eqs. (1) and (2), the heavy meson fields of  $\bar{Q}q$  are parametrized by the heavy pseudoscalar and vector mesons,

$$H_a = \frac{1 + \not{v}}{2} [P_{a\mu}^* \gamma^\mu - P_a \gamma_5] , \quad (3)$$

$$\bar{H}_a = \gamma_0 H_a^\dagger \gamma_0 , \quad (4)$$

where the subscripts  $a, b$  are for light flavors ( $u, d$ ). The pseudoscalar and vector fields are

normalized as

$$\langle 0|P|P(p_\mu)\rangle = \sqrt{p^0}, \quad (5)$$

$$\langle 0|P_\mu^*|P^*(p_\mu, \lambda)\rangle = \epsilon(\lambda)_\mu \sqrt{p^0}, \quad (6)$$

where  $\epsilon(\lambda)_\mu$  is the polarization vector of  $P^*$  with polarization  $\lambda$ . The axial and vector currents of light flavors are given by

$$A^\mu = \frac{1}{2}(\xi^\dagger \partial^\mu \xi - \xi \partial^\mu \xi^\dagger), \quad (7)$$

$$V^\mu = \frac{1}{2}(\xi^\dagger \partial^\mu \xi + \xi \partial^\mu \xi^\dagger), \quad (8)$$

where  $\xi = \exp(i\hat{\pi}/f_\pi)$ ,  $f_\pi = 132$  MeV is the pion decay constant and we define the pion field by

$$\hat{\pi} = \begin{pmatrix} \frac{\pi^0}{\sqrt{2}} & \pi^+ \\ \pi^- & -\frac{\pi^0}{\sqrt{2}} \end{pmatrix}. \quad (9)$$

Finally the vector ( $\rho$  and  $\omega$ ) meson field and its field tensor are defined by

$$\rho_\mu = i \frac{g_V}{\sqrt{2}} \hat{\rho}_\mu, \quad (10)$$

$$\hat{\rho}_\mu = \begin{pmatrix} \frac{\rho^0}{\sqrt{2}} + \frac{\omega}{\sqrt{2}} & \rho^+ \\ \rho^- & -\frac{\rho^0}{\sqrt{2}} + \frac{\omega}{\sqrt{2}} \end{pmatrix}_\mu, \quad (11)$$

$$F_{\mu\nu}(\rho) = \partial_\mu \rho_\nu - \partial_\nu \rho_\mu + [\rho_\mu, \rho_\nu], \quad (12)$$

where  $g_V$  is the gauge coupling constant of the hidden local symmetry [50].

The coupling constant  $g_\pi$  for  $\pi PP^*$  is fixed from the strong decay of  $D^* \rightarrow D\pi$  [51]. The coupling constants  $\beta$  and  $\lambda$  are determined by the radiative decays of  $D^*$  and semileptonic decays of  $B$  with the vector meson dominance [52]. The resulting values are given in Table. II.

From Eq. (1) we obtain  $\pi PP^*$  and  $\pi P^* P^*$  vertices as

$$\mathcal{L}_{\pi PP^*} = 2 \frac{g_\pi}{f_\pi} (P_a^\dagger P_{b\mu}^* + P_{a\mu}^* P_b) \partial^\mu \hat{\pi}_{ab}, \quad (13)$$

$$\mathcal{L}_{\pi P^* P^*} = 2i \frac{g_\pi}{f_\pi} \epsilon^{\mu\nu\alpha\beta} v_\mu P_{a\beta}^* P_{b\nu}^* \partial_\alpha \hat{\pi}_{ab}. \quad (14)$$

We note that there is no  $\pi PP$  vertex due to parity invariance. Similarly, from Eq. (2), we

TABLE II: Masses and coupling constants of mesons.

	$m_\alpha$ [MeV]	$g_\pi$	$\beta$	$\lambda$ [GeV $^{-1}$ ]	$g_{\alpha NN}^2/4\pi$	$\kappa$
$\pi$	137.27	0.59	—	—	13.6	—
$\rho$	769.9	—	0.9	0.56	0.84	6.1
$\omega$	781.94	—	0.9	0.56	20.0	0.0

derive the vector meson vertices as

$$\mathcal{L}_{vPP} = -\sqrt{2}\beta g_V P_b P_a^\dagger v \cdot \hat{\rho}_{ba}, \quad (15)$$

$$\mathcal{L}_{vPP^*} = -2\sqrt{2}\lambda g_V v_\mu \epsilon^{\mu\nu\alpha\beta} \left( P_a^\dagger P_{b\beta}^* - P_{a\beta}^{*\dagger} P_b \right) \partial_\nu (\hat{\rho}_\alpha)_{ba}, \quad (16)$$

$$\begin{aligned} \mathcal{L}_{vP^*P^*} = & \sqrt{2}\beta g_V P_b^* P_a^{*\dagger} v \cdot \hat{\rho}_{ba} \\ & + i2\sqrt{2}\lambda g_V P_{a\mu}^{*\dagger} P_{b\nu}^* (\partial^\mu (\hat{\rho}^\nu)_{ba} - \partial^\nu (\hat{\rho}^\mu)_{ba}). \end{aligned} \quad (17)$$

The interaction Lagrangians for a meson and nucleons are given by the standard form,

$$\mathcal{L}_{\pi NN} = \sqrt{2}i g_{\pi NN} \bar{N} \gamma_5 \hat{\pi} N, \quad (18)$$

$$\mathcal{L}_{vNN} = \sqrt{2} g_{vNN} \left[ \bar{N} \gamma_\mu \hat{\rho}^\mu N + \frac{\kappa}{2m_N} \bar{N} \sigma_{\mu\nu} \partial^\nu \hat{\rho}^\mu N \right], \quad (19)$$

where  $N = (p, n)^T$  is the nucleon field. In the vector meson interaction there are vector (Dirac) and tensor (Pauli) terms. For  $\rho$  meson the tensor term dominates, while for  $\omega$  it is negligible. The coupling constants associated to the nucleon are taken from Ref. [53] as summarized also in Table. II.

To parametrize internal structure of hadrons, we introduce form factors associated with finite size of the mesons and nucleons. We adopt a dipole form factor at each vertex :

$$F_\alpha(\Lambda, \vec{q}) = \frac{\Lambda^2 - m_\alpha^2}{\Lambda^2 + |\vec{q}|^2}, \quad (20)$$

where  $m_\alpha$  and  $\vec{q}$  are the mass and three-momentum of the incoming meson  $\alpha$  ( $= \pi, \rho, \omega$ ). The cutoff parameter for the meson-nucleon vertex  $\Lambda = \Lambda_N$  is determined such that the resulting  $NN$  potential reproduces the binding energy of the deuteron. For the derivation of the potential, see the discussion below and also Appendix. When the  $NN$  potential is constructed only by  $\pi$  exchange,  $\Lambda_N = 830$  MeV, while when the potential is constructed by  $\pi, \rho, \omega$  exchanges,  $\Lambda_N = 846$  MeV. To test the validity of the potentials, we have computed

low energy properties of the deuteron and  $NN$  scattering. The results are given in Table. III, for the  $\pi$  exchange and  $\pi, \rho, \omega$  exchange potentials. Another cutoff parameter  $\Lambda_P$  for the meson-meson vertex is determined by the ratio of the size of the pseudoscalar meson  $P$  as given in Ref. [16],  $\Lambda_D = 1.35\Lambda_N$  and  $\Lambda_B = 1.29\Lambda_N$ . We have adopted the same value of the cutoff for the vertices including vector meson  $P^*$ .

We have also constructed the interaction from the flavor  $SU(4)$  symmetry for comparison with that of the heavy quark symmetry. However, it turns out that the coupling strength of the  $\pi PP^*$  vertex is smaller than the one determined from the experimental decay rate of  $D^* \rightarrow D\pi$ . We have found that there is no bound state when such a small coupling strength is used. In this work we use the potential derived from the heavy quark symmetry with the coupling strength determined by the experimental decay width of  $D^* \rightarrow D\pi$ .

Having all the above interaction Lagrangians for the Yukawa vertices, we obtain potentials for various channels which are summarized in Appendix. In doing so, we employ the static approximation where the energy transfer can be ignored. This is a good approximation when  $P^*$  does not decay into  $P\pi$  in the heavy quark limit and at low energies below the threshold of pion productions. The total Hamiltonian is then given as the sum of the kinetic energies and the potential for coupled channels as shown explicitly in Appendix. We solve the coupled Schrödinger equations for the  $PN$  and  $P^*N$  system in their center of mass frame by using the numerical method developed in Ref. [54]. The total energy  $E$  is measured from the  $PN$  threshold.

We have tested our method for the deuteron case where we can compare our results with the known results. As summarized in Table. III, we point out that the  $\pi$  exchange potential reproduces the deuteron properties well, such as the binding energy  $E_B$ , the  $D$ -wave probability  $P_D$ , the mean square radius  $\langle r^2 \rangle^{1/2}$ , as well as the scattering length  $a$  and the effective range  $r_e$  in the  $^3S_1$  and  $I = 0$  channel in the  $NN$  scattering. The results with the  $\pi, \rho, \omega$  exchange potential are very similar to those of the  $\pi$  exchange potential. This indicates that the  $\rho$  and  $\omega$  exchanges play only a minor role for low energy properties as expected, because the deuteron is a loosely bound state and a rather extended object.

TABLE III: Low energy properties of the  $NN$  system. Results for the  $\pi$  and  $\pi\rho\omega$  potentials are compared.

Potential	$\Lambda_N$ [MeV]	$E_B$ [MeV]	$P_D$ [%]	$\langle r^2 \rangle^{1/2}$ [fm]	$a$ [fm]	$r_e$ [fm]
$\pi$	830	2.22	5.4	3.7	5.27	1.50
$\pi\rho\omega$	846	2.22	5.3	3.7	5.23	1.49

TABLE IV: Binding energies, root mean square radii and cutoff parameters of heavy mesons. Results for the  $\pi$  and  $\pi\rho\omega$  potentials are compared.

	$\bar{D}N(\pi)$	$\bar{D}N(\pi\rho\omega)$	$BN(\pi)$	$BN(\pi\rho\omega)$
$E_B$ [MeV]	1.60	2.14	19.50	23.04
$\langle r^2 \rangle^{1/2}$ [fm]	3.5	3.2	1.3	1.2
$\Lambda_P$ [MeV]	1121	1142	1070	1091

### III. BOUND STATES

In this section let us study low lying states for the  $PN$ - $P^*N$  system. We find a bound state in the  $(I, J^P) = (0, 1/2^-)$  channel as discussed in Ref. [16]. The results of  $\pi$  exchange potential are almost the same as those of  $\pi, \rho, \omega$  exchange potential as seen in the deuteron case. The binding energies  $E_B = |E|$  and the root mean square radii  $\langle r^2 \rangle^{1/2}$  for  $\bar{D}N$  and  $BN$  states are shown in Table. IV. As emphasized in Ref. [16], the tensor force causing mixing between  $\bar{D}N$  and  $\bar{D}^*N$  states of different angular momenta by  $\Delta L = 2$  yields strong attraction. This is particularly so for heavier quark sector, where  $P$  and  $P^*$  mesons degenerate. Therefore, the  $BN$  state is more bound than the  $\bar{D}N$  state. In fact, the binding energies of  $BN$  states are about ten times larger than those of  $\bar{D}N$  states. The mass dependence of the binding energies will be discussed more in detail in section V.

The corresponding root mean square radii are over 3 fm for the charm and over 1 fm for the bottom baryons, respectively. Both radii are larger than typical hadron size of order 1 fm, justifying the hadronic composite structure of the present states. To complete our presentations, in Fig. 1, we also show bound state wave functions when the  $\pi\rho\omega$  potential is used. Finally we note that there is no bound state in  $J^P = 3/2^-$  as shown in Ref. [16]



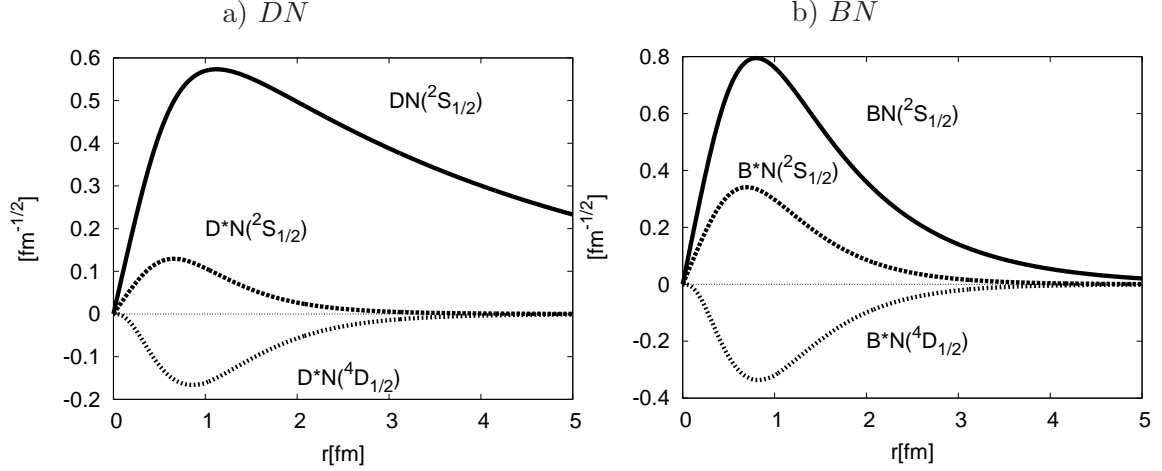


FIG. 1: The wave functions of the  $\bar{D}N$  and  $BN$  bound states with  $(I, J^P) = (0, 1/2^-)$  when the  $\pi\rho\omega$  potential is used.

#### IV. SCATTERING STATES AND RESONANCE

Let us now turn to the scattering region above the  $PN$  threshold. First, we see the scattering state in the  $(I, J^P) = (0, 1/2^-)$  channel. In Fig. 2, we present the phase shifts  $\delta$  of  $PN(^2S_{1/2})$ ,  $P^*N(^2S_{1/2})$  and  $P^*N(^4D_{1/2})$  channels when the  $\pi\rho\omega$  potential is used. Each phase shift is plotted as a function of the scattering energy  $E$  in the center of mass system. The  $PN$  phase shift starts at  $\delta = \pi$  because of the presence of the bound state discussed in the previous section. Otherwise the energy dependence of all the phase shifts is rather smooth. We summarize the scattering lengths and the effective ranges in Table V. We have checked the relation  $E_B = \frac{1}{2\mu a^2}$  ( $\mu$  is reduced mass) holds with good accuracy within a few percents. We also note that the properties of  $\bar{D}N$  system are rather similar to those of the  $NN$  system.

We show the total cross sections of  $\bar{D}N$  and  $BN$  scattering when the  $\pi\rho\omega$  potential is used in Fig. 3. They start from the maximum value at the threshold and decrease monotonically. For shallower bound state (for  $\bar{D}N$  system), the peak value at the threshold is larger, due to the presence of the bound state near the threshold. In the limit the binding energy  $E_B \rightarrow 0$ , the peak diverges as explained by the zero-energy resonance.

In the  $(I, J^P) = (0, 3/2^-)$  channel, we find an interesting structure; the phase shifts shown in Fig. 4 indicate a resonance at the scattering energy  $E_{re} = 113.19$  MeV for  $\bar{D}N$  and at  $E_{re} = 6.93$  MeV for  $BN$ , as the phase shifts cross  $\pi/2$ . The mechanism of the resonance

TABLE V: The scattering length  $a$  and effective range  $r_e$  with  $(I, J^P) = (0, 1/2^-)$ .

	$\bar{D}N(\pi)$	$\bar{D}N(\pi\rho\omega)$	$BN(\pi)$	$BN(\pi\rho\omega)$
$a$ [fm]	4.36	4.38	1.61	1.56
$r_e$ [fm]	1.04	1.05	0.71	0.68

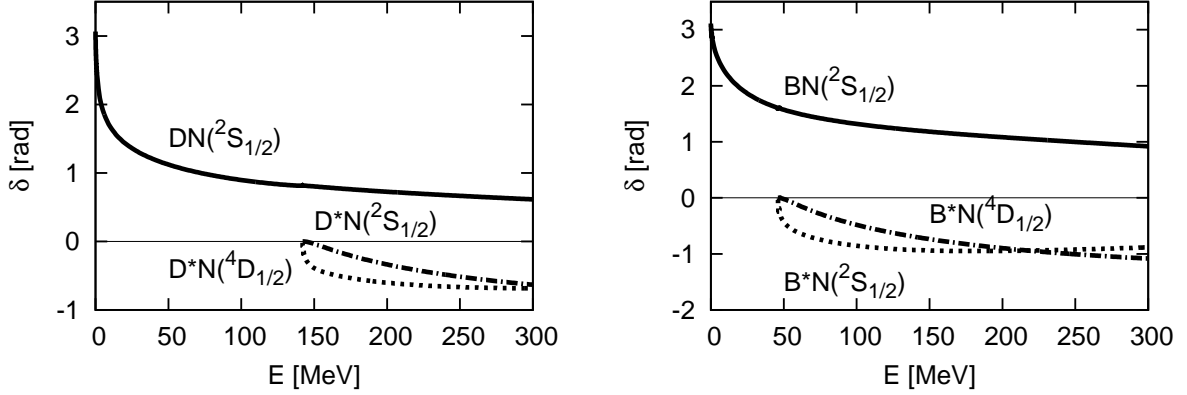


FIG. 2: Phase shift of the  $\bar{D}N$  and  $BN$  scattering state with  $(I, J^P) = (0, 1/2^-)$  when the  $\pi\rho\omega$  potential is used.

can be understood by looking at the phase shifts in other channels; in particular, we find that those of  $\bar{D}^*N(^4S_{3/2})$  and of  $B^*N(^4S_{3/2})$  start from  $\delta = \pi$ , indicating the presence of a bound state in these channels. Indeed, we checked that, when the  $PN(^2D_{3/2})$  channel is ignored and only the  $P^*N(^2D_{3/2})$ ,  $P^*N(^4D_{3/2})$  and  $P^*N(^4S_{3/2})$  channels are considered, there are bound states at  $E_B = 11.50$  MeV from the  $\bar{D}^*N$  threshold and at  $E_B = 21.67$  MeV from the  $B^*N$  threshold. Therefore, these resonances are the Feshbach resonances.

We can also extract the width of the resonances from the slope of the phase shift at  $\pi/2$ ;

$$\Gamma = \frac{2}{d\delta/dE|_{E=E_{re}}}, \quad (21)$$

for the phase shift  $\delta$  in the partial wave, namely  $PN(^2D_{3/2})$ . The results are listed in Table VI. These values are very small in particular for the  $BN$  system. There are two reasons for that; one is that the resonance energy is close to the threshold of the open channel of  $BN$ , and the other is that the decay occurs in the  $D$ -wave. The latter effect is important as the decay width is proportional to the fifth power ( $2L + 1$ ,  $L = 2$ ) of the

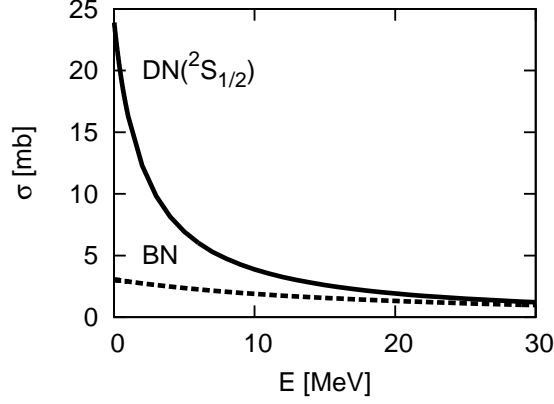


FIG. 3: Total cross section of the  $\bar{D}N$  and  $BN$  scattering state with  $(I, J^P) = (0, 1/2^-)$  when the  $\pi\rho\omega$  potential is used.

decaying momentum and works as a suppression factor near the threshold. The physics behind is the centrifugal barrier for a partial wave with a finite angular momentum.

In Fig. 5, we plot the cross section in the  $(I, J^P) = (0, 3/2^-)$  channel. Because there is a resonant state at  $E_{re} = 113$  MeV in the  $\bar{D}N$  scattering and at  $E_{re} = 7$  MeV in the  $BN$  scattering, the cross section becomes maximum at each resonance energy.

TABLE VI: The resonance energy and decay width for  $(I, J^P) = (0, 3/2^-)$ .

	$\bar{D}N(\pi)$	$\bar{D}N(\pi\rho\omega)$	$BN(\pi)$	$BN(\pi\rho\omega)$
$E_{re}$ [MeV]	113.51	113.19	8.41	6.93
$\Gamma$ [MeV]	19.43	17.72	0.16	0.0946

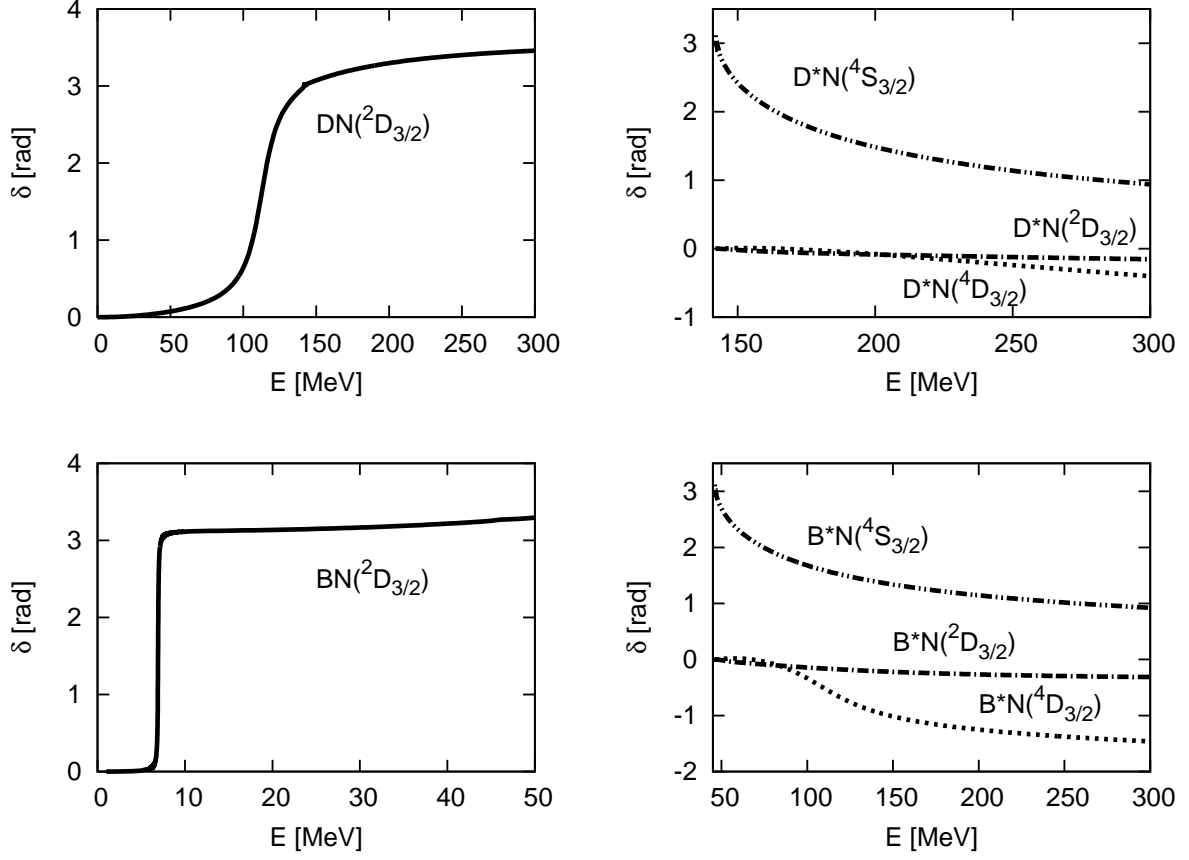


FIG. 4: Phase shift of the  $\bar{D}N$  and  $BN$  scattering state with  $(I, J^P) = (0, 3/2^-)$  when the  $\pi\rho\omega$  potential is used.

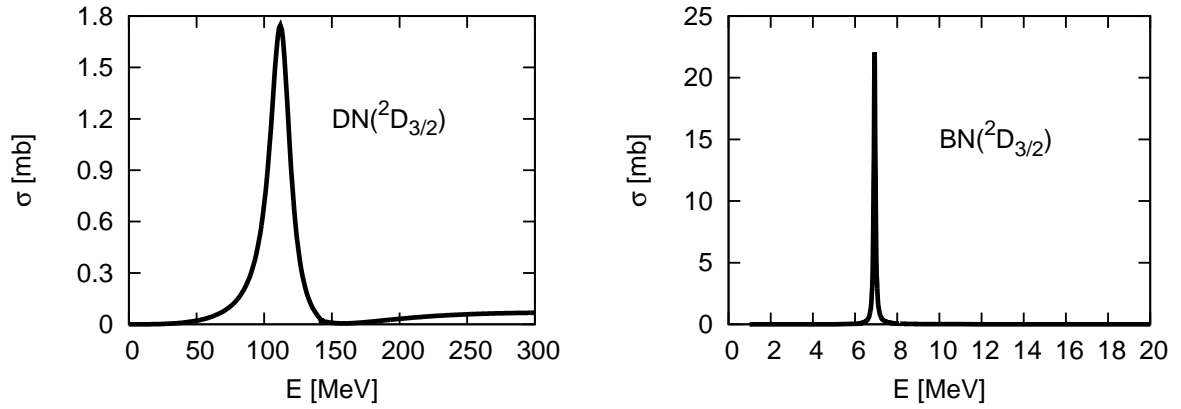


FIG. 5: Total cross section of the  $\bar{D}N$  and  $BN$  scattering state with  $(I, J^P) = (0, 3/2^-)$  when the  $\pi\rho\omega$  potential is used.

## V. FLAVOR DEPENDENCE OF THE BOUND AND RESONANT STATE

In the previous sections, we have seen that the bound and resonant states in the bottom sector receive more attractions than in the charm sector. In other words, the existence of the present exotic baryons is a unique feature of heavier mesons. To see this more quantitatively, let us vary the mass of the heavy mesons continuously and see how the bound and resonant state properties change. Let us make an interpolation of heavy meson masses, by regarding the mass of the  $P$  meson as a function of the mass of the  $P^*$  meson. Given the experimental values of  $(m_K, m_{K^*})$ ,  $(m_{\bar{D}}, m_{\bar{D}^*})$  and  $(m_B, m_{B^*})$ , we find the parametrization  $m_{P^*} - m_P = 1.9 \times 10^6 / m_{P^*}^{1.25}$  in units of MeV to optimally interpolate the masses at the three flavor points, as shown in Fig. 6. In the heavy quark mass limit, the mass difference  $m_{P^*} - m_P$  is expected to be proportional to the inverse mass of the heavy quark and so to the mass of the vector meson. The deviation of the power 1.25 from unity is due to some finite mass corrections. Here, the details of the functional form is not important, but only an interpolation at a qualitative level is enough.

Having this parametrization, we perform calculations for the bound and resonant state properties. In Fig. 7, the eigenenergy of the bound state is plotted as a function of  $m_{P^*}$ . As expected, the binding energy increases as the mass of the heavy meson increases. What is then interesting is that the bound state disappears at  $m_{P^*} \sim 1700$  MeV, which is about 300 MeV smaller than the mass of  $D^*$ . Hence, that the masses of the mesons are heavy is crucial for the presence of the exotic baryons.

In Fig. 8 we plot the resonance energy as a function of  $m_{P^*}$ ; it increases as  $m_{P^*}$  decreases. The resonance exists for a small meson mass region also where the bound state no longer exists. The behavior of the decay width of the resonance as shown in Fig. 9 is interesting, as it takes the maximum value at  $m_{P^*} \sim 1700$  MeV. For larger masses  $m_{P^*} \sim 5400$  MeV and beyond, the width becomes zero, where the would-be bound state of the single channel  $P^*N$  is located below the decay channel of  $PN$ . Contrary, as  $m_{P^*}$  decreases, the resonance energy becomes larger and its wave function extends more. This suppresses the overlap of the wave functions of the resonance and decaying channels. These explain the reason that the decay width takes the maximum value at a medium energy point.

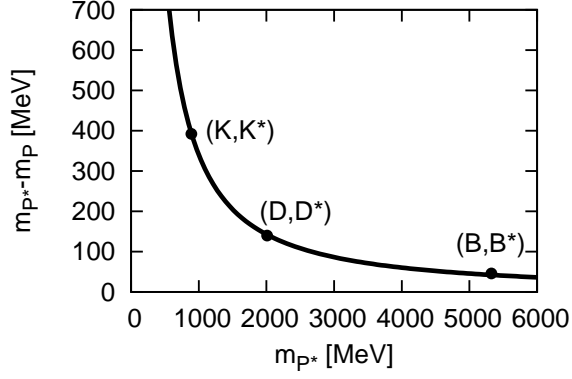


FIG. 6: The mass splitting  $m_{P^*} - m_P$  as a function of the heavy vector meson mass  $m_{P^*}$ .

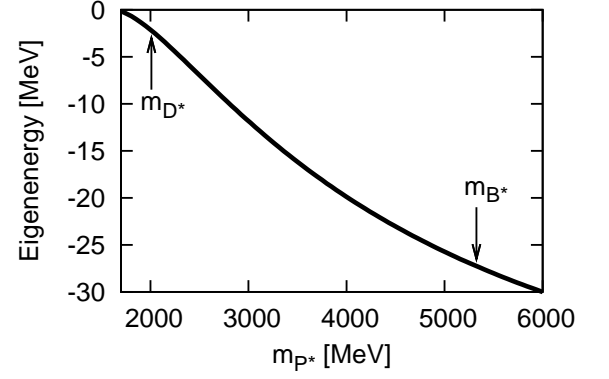


FIG. 7: Eigenenergy for  $(I, J^P) = (0, 1/2^-)$  as a function of heavy vector meson mass  $m_{P^*}$ .

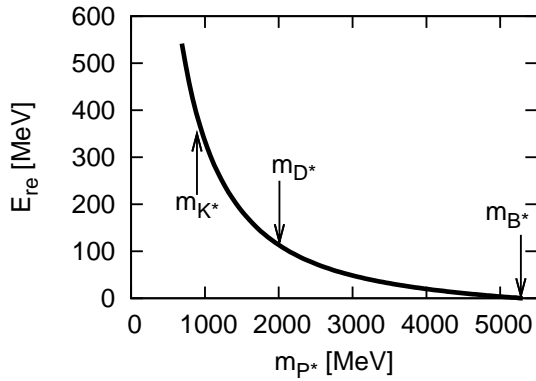


FIG. 8: Resonance energy for  $(I, J^P) = (0, 3/2^-)$  as a function of heavy vector meson mass  $m_{P^*}$ .

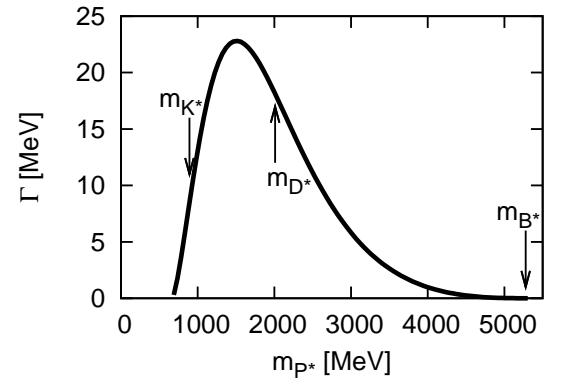


FIG. 9: Decay width for  $(I, J^P) = (0, 3/2^-)$  as a function of heavy vector meson mass  $m_{P^*}$ .

## VI. SUMMARY

In this paper, we have investigated heavy baryons having exotic flavor quantum numbers as hadronic composites ( $\bar{D}N$  and  $BN$ ) of a heavy meson and a nucleon. In the quark content, this has minimally five quarks of  $qqqq\bar{Q}$ . Because of the light flavor content, the one pion exchange interaction is at work and its tensor force plays a very important role through channel couplings of different angular momentum states of  $\Delta L = 2$ . The mechanism is essentially the same as that for the deuteron binding. Such channel coupling effects become more striking if the masses of the mesons become heavier than those of charmed mesons, where the pseudoscalar and vector mesons are more degenerate.

Our interactions were determined consistently with the low energy properties of the two-nucleon systems including the deuteron and some properties of the heavy mesons and the nucleon. We have confirmed that the previous finding of the pion dominance for such novel states to exist by including short range interaction mediated by  $\rho$  and  $\omega$ . As a result, we have found that the bound state appears for  $(I, J^P) = (0, 1/2^-)$ , which has been predicted before. We have also found a new resonant state for  $(I, J^P) = (0, 3/2^-)$  with a narrow decay width as a Feshbach resonance predominated by the heavy vector meson and nucleon bound state, which decays into an open channel of a pseudoscalar and a nucleon. We have investigated the flavor dependence by changing continuously the masses of heavy mesons. Then we have found that the bound states in the  $(I, J^P) = (0, 1/2^-)$  channel exist for the vector meson masses larger than 1700 MeV. We have also found a resonance with  $(I, J^P) = (0, 3/2^-)$  having a narrow decay width. Therefore the heavy masses of the mesons are crucial for the existence of the exotic baryons in the present study.

Final remark is related to how these states are observed. Resonant states may be found as a pair of strongly decaying pseudoscalar meson and nucleon in high energy hadronic collisions. Exclusive experiments such as hadron and photon induced reactions would be useful for the observation of bound states since they are stable against strong decay. The understanding of production reactions is an important issue in the future, which can be studied at J-PARC, GSI and so on. It has been recently proposed that the quark-gluon plasma formed in the relativistic heavy ion collisions will serve a source of multi-particles including exotic hadrons [55]. Therefore, RHIC, LHC and other facilities will also help to search the exotic states predicted in the present work.

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## Appendix A: Potentials and kinetic terms

Interaction potentials are derived by using the Lagrangians Eqs. (13)-(19). In deriving the potentials we use the static approximation where the energy transfer can be neglected as compared to the momentum transfer. The resulting potentials for the coupled channel systems are given in the matrix form of  $3 \times 3$  for  $J^P = 1/2^-$  and of  $4 \times 4$  for  $J^P = 3/2^-$ ,

$$V_{1/2^-} = \begin{pmatrix} V_{1/2^-}^{11} & V_{1/2^-}^{12} & V_{1/2^-}^{13} \\ V_{1/2^-}^{21} & V_{1/2^-}^{22} & V_{1/2^-}^{23} \\ V_{1/2^-}^{31} & V_{1/2^-}^{32} & V_{1/2^-}^{33} \end{pmatrix}, \quad (\text{A1})$$

$$V_{3/2^-} = \begin{pmatrix} V_{3/2^-}^{11} & V_{3/2^-}^{12} & V_{3/2^-}^{13} & V_{3/2^-}^{14} \\ V_{3/2^-}^{21} & V_{3/2^-}^{22} & V_{3/2^-}^{23} & V_{3/2^-}^{24} \\ V_{3/2^-}^{31} & V_{3/2^-}^{32} & V_{3/2^-}^{33} & V_{3/2^-}^{34} \\ V_{3/2^-}^{41} & V_{3/2^-}^{42} & V_{3/2^-}^{43} & V_{3/2^-}^{44} \end{pmatrix}, \quad (\text{A2})$$

with the basis given in Table I in the same ordering. The  $\pi$  exchange potential between heavy flavor meson and nucleon is obtained by

$$V_{1/2^-}^\pi = \frac{g_\pi g_{\pi NN}}{\sqrt{2}m_N f_\pi} \frac{1}{3} \begin{pmatrix} 0 & \sqrt{3}C_{m_\pi} & -\sqrt{6}T_{m_\pi} \\ \sqrt{3}C_{m_\pi} & -2C_{m_\pi} & -\sqrt{2}T_{m_\pi} \\ -\sqrt{6}T_{m_\pi} & -\sqrt{2}T_{m_\pi} & C_{m_\pi} - 2T_{m_\pi} \end{pmatrix} \vec{\tau}_P \cdot \vec{\tau}_N, \quad (\text{A3})$$

$$V_{3/2^-}^\pi = \frac{g_\pi g_{\pi NN}}{\sqrt{2}m_N f_\pi} \frac{1}{3} \begin{pmatrix} 0 & \sqrt{3}T_{m_\pi} & -\sqrt{3}T_{m_\pi} & \sqrt{3}C_{m_\pi} \\ \sqrt{3}T_{m_\pi} & C_{m_\pi} & 2T_{m_\pi} & T_{m_\pi} \\ -\sqrt{3}T_{m_\pi} & 2T_{m_\pi} & C_{m_\pi} & -T_{m_\pi} \\ \sqrt{3}C_{m_\pi} & T_{m_\pi} & -T_{m_\pi} & -2C_{m_\pi} \end{pmatrix} \vec{\tau}_P \cdot \vec{\tau}_N, \quad (\text{A4})$$



where  $C_m = C(r; m)$ ,  $T_m = T(r; m)$ , and  $\vec{\tau}_P$  and  $\vec{\tau}_N$  are the isospin matrices for  $P(P^*)$  and  $N$ . The functions  $C(r; m)$  and  $T(r; m)$  are given by

$$C(r; m) = \int \frac{d^3q}{(2\pi)^3} \frac{m^2}{\vec{q}^2 + m^2} e^{i\vec{q}\cdot\vec{r}} F(\Lambda_P, \vec{q}) F(\Lambda_N, \vec{q}), \quad (\text{A5})$$

$$T(r; m) S_{12}(\hat{r}) = \int \frac{d^3q}{(2\pi)^3} \frac{-\vec{q}^2}{\vec{q}^2 + m^2} S_{12}(\hat{q}) e^{i\vec{q}\cdot\vec{r}} F(\Lambda_P, \vec{q}) F(\Lambda_N, \vec{q}), \quad (\text{A6})$$

with  $S_{12}(\hat{x}) = 3(\vec{\sigma}_1 \cdot \hat{x})(\vec{\sigma}_2 \cdot \hat{x}) - \vec{\sigma}_1 \cdot \vec{\sigma}_2$ , and  $F(\Lambda, \vec{q})$  denotes the form factor given in Eq. (20).

The corresponding potentials of the  $\rho$  meson exchange are given by

$$\begin{aligned} V_{1/2-}^\rho = & \frac{g_V g_{\rho NN} \beta}{\sqrt{2} m_\rho^2} \begin{pmatrix} C_{m_\rho} & 0 & 0 \\ 0 & C_{m_\rho} & 0 \\ 0 & 0 & C_{m_\rho} \end{pmatrix} \vec{\tau}_P \cdot \vec{\tau}_N \\ & + \frac{g_V g_{\rho NN} \lambda(1 + \kappa)}{\sqrt{2} m_N} \frac{1}{3} \begin{pmatrix} 0 & 2\sqrt{3}C_{m_\rho} & \sqrt{6}T_{m_\rho} \\ 2\sqrt{3}C_{m_\rho} & -4C_{m_\rho} & \sqrt{2}T_{m_\rho} \\ \sqrt{6}T_{m_\rho} & \sqrt{2}T_{m_\rho} & 2C_{m_\rho} + 2T_{m_\rho} \end{pmatrix} \vec{\tau}_P \cdot \vec{\tau}_N, \end{aligned} \quad (\text{A7})$$

$$\begin{aligned} V_{3/2-}^\rho = & \frac{g_V g_{\rho NN} \beta}{\sqrt{2} m_\rho^2} \begin{pmatrix} C_{m_\rho} & 0 & 0 & 0 \\ 0 & C_{m_\rho} & 0 & 0 \\ 0 & 0 & C_{m_\rho} & 0 \\ 0 & 0 & 0 & C_{m_\rho} \end{pmatrix} \vec{\tau}_P \cdot \vec{\tau}_N \\ & + \frac{g_V g_{\rho NN} \lambda(1 + \kappa)}{\sqrt{2} m_N} \frac{1}{3} \begin{pmatrix} 0 & -\sqrt{3}T_{m_\rho} & \sqrt{3}T_{m_\rho} & 2\sqrt{3}C_{m_\rho} \\ -\sqrt{3}T_{m_\rho} & 2C_{m_\rho} & -2T_{m_\rho} & -T_{m_\rho} \\ \sqrt{3}T_{m_\rho} & -2T_{m_\rho} & 2C_{m_\rho} & T_{m_\rho} \\ 2\sqrt{3}C_{m_\rho} & -T_{m_\rho} & T_{m_\rho} & -4C_{m_\rho} \end{pmatrix} \vec{\tau}_P \cdot \vec{\tau}_N. \end{aligned} \quad (\text{A8})$$

The  $\omega$  meson exchange potential can be obtained by replacing the relevant coupling constants and the mass of the exchanged meson, and by removing the isospin factor  $\vec{\tau}_P \cdot \vec{\tau}_N$ . The anomalous coupling  $\kappa$  for the  $\omega$  meson exchange potential is set as zero in Eqs. (A7)-(A8).

In Figs. 10-12 the functional forms of all the potentials are shown. In these figures, we can see clearly the dominance of the tensor force in the transition amplitude, for instance,  $V_{1/2-}^{13}$  and  $V_{1/2-}^{23}$ .

The kinetic terms are given by

$$K_{1/2^-} = \text{diag} \left( -\frac{1}{2\tilde{m}_P} \Delta_0, -\frac{1}{2\tilde{m}_{P^*}} \Delta_0 + \Delta m_{PP^*}, -\frac{1}{2\tilde{m}_{P^*}} \Delta_2 + \Delta m_{PP^*} \right), \quad (\text{A9})$$

$$K_{3/2^-} = \text{diag} \left( -\frac{1}{2\tilde{m}_P} \Delta_2, -\frac{1}{2\tilde{m}_{P^*}} \Delta_0 + \Delta m_{PP^*}, -\frac{1}{2\tilde{m}_{P^*}} \Delta_2 + \Delta m_{PP^*}, \right. \\ \left. -\frac{1}{2\tilde{m}_{P^*}} \Delta_2 + \Delta m_{PP^*} \right), \quad (\text{A10})$$

for  $J^P = 1/2^-$  and  $3/2^-$ , respectively. Here, we define  $\Delta_0 = \partial^2/\partial r^2 + (2/r)\partial/\partial r$  and  $\Delta_2 = \Delta_0 + 6/r^2$ ,  $\tilde{m}_{P^{(*)}} = m_N m_{P^{(*)}}/(m_N + m_{P^{(*)}})$ , with  $\Delta m_{PP^*} = m_{P^*} - m_P$ . The total Hamiltonian is then given by  $H_{J^P} = K_{J^P} + V_{J^P}$ .

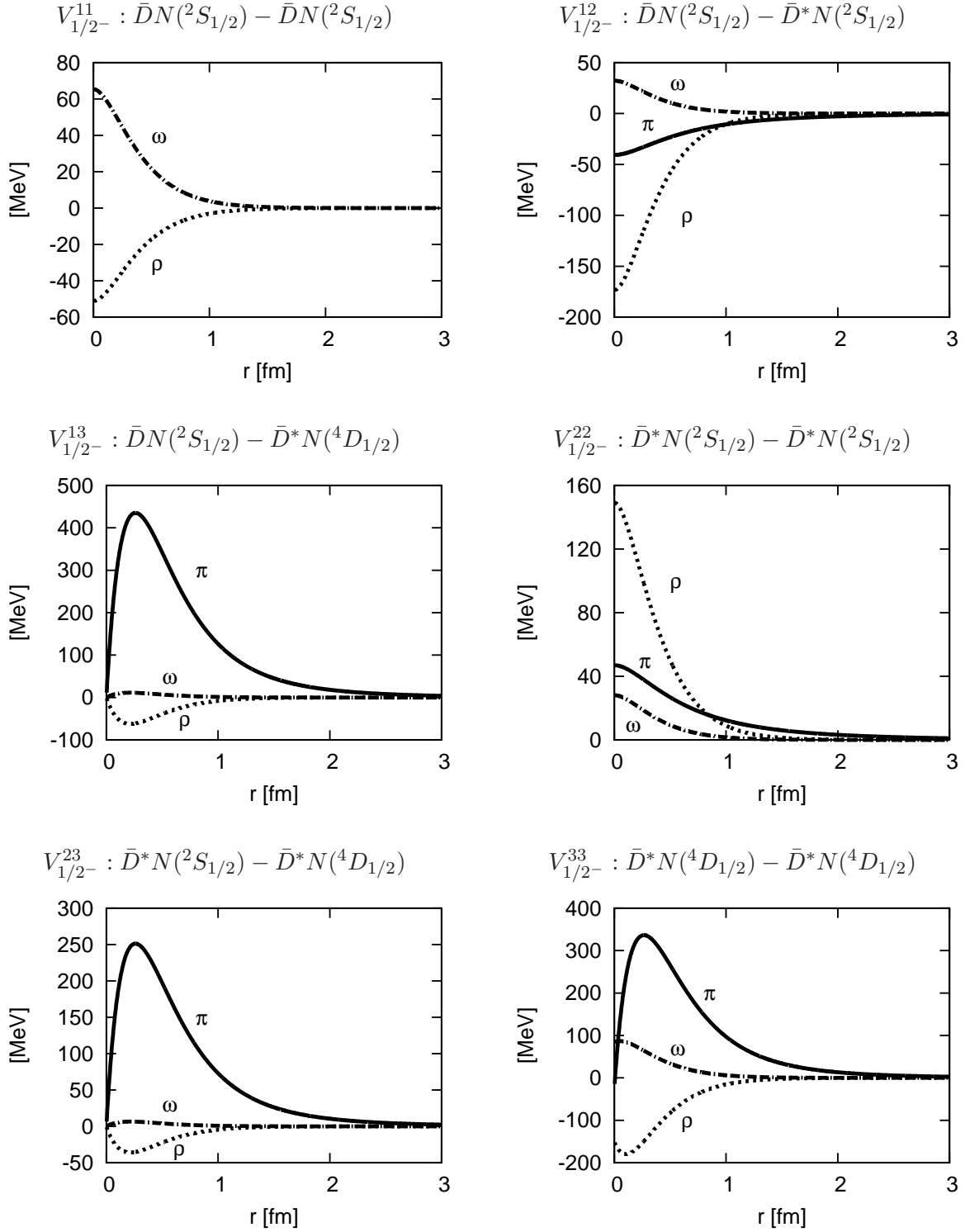


FIG. 10: Various components of the  $\pi\rho\omega$  exchange potential for  $(I, J^P) = (0, 1/2^-)$ .

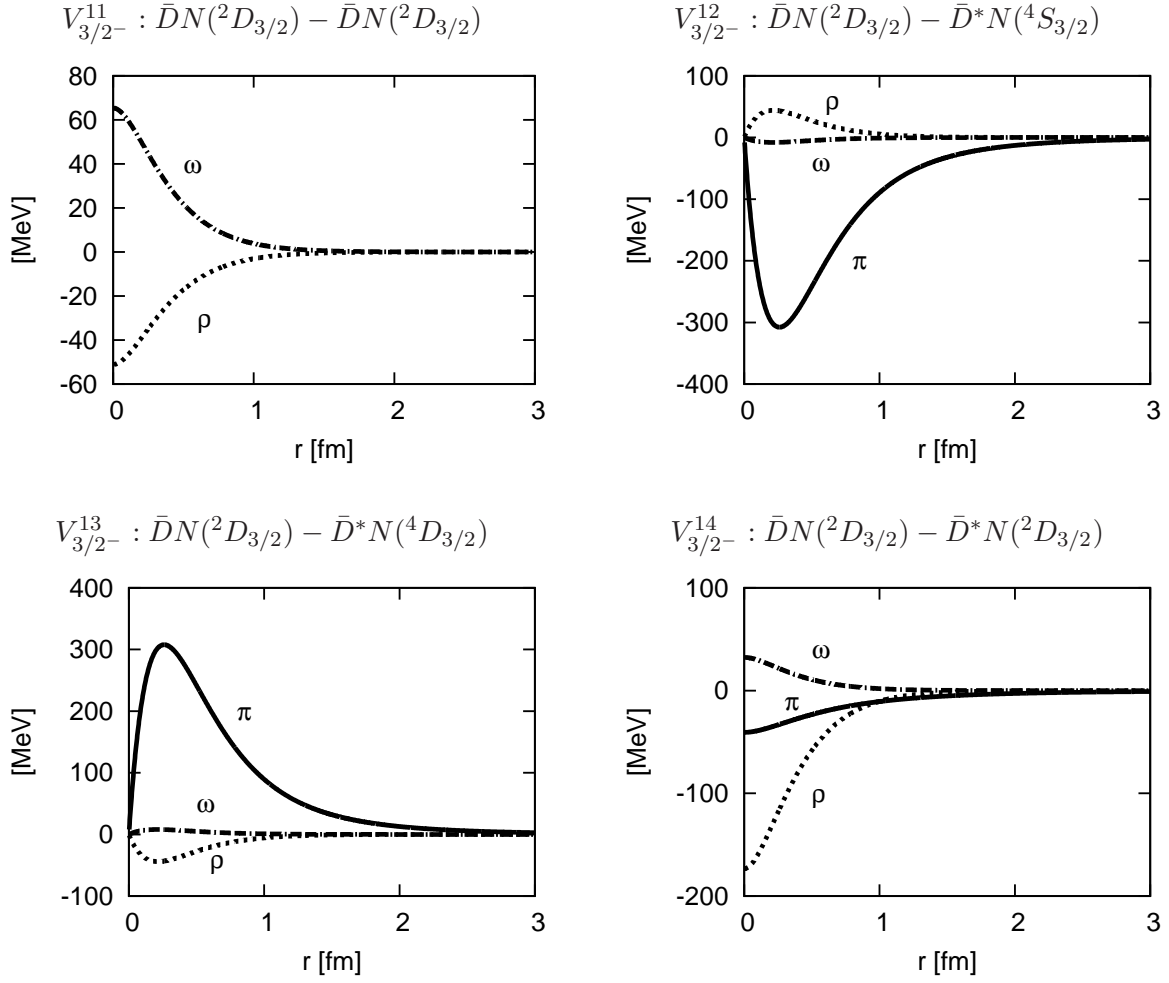


FIG. 11: Various components of the  $\pi\rho\omega$  exchange potential for  $(I, J^P) = (0, 3/2^-)$ .

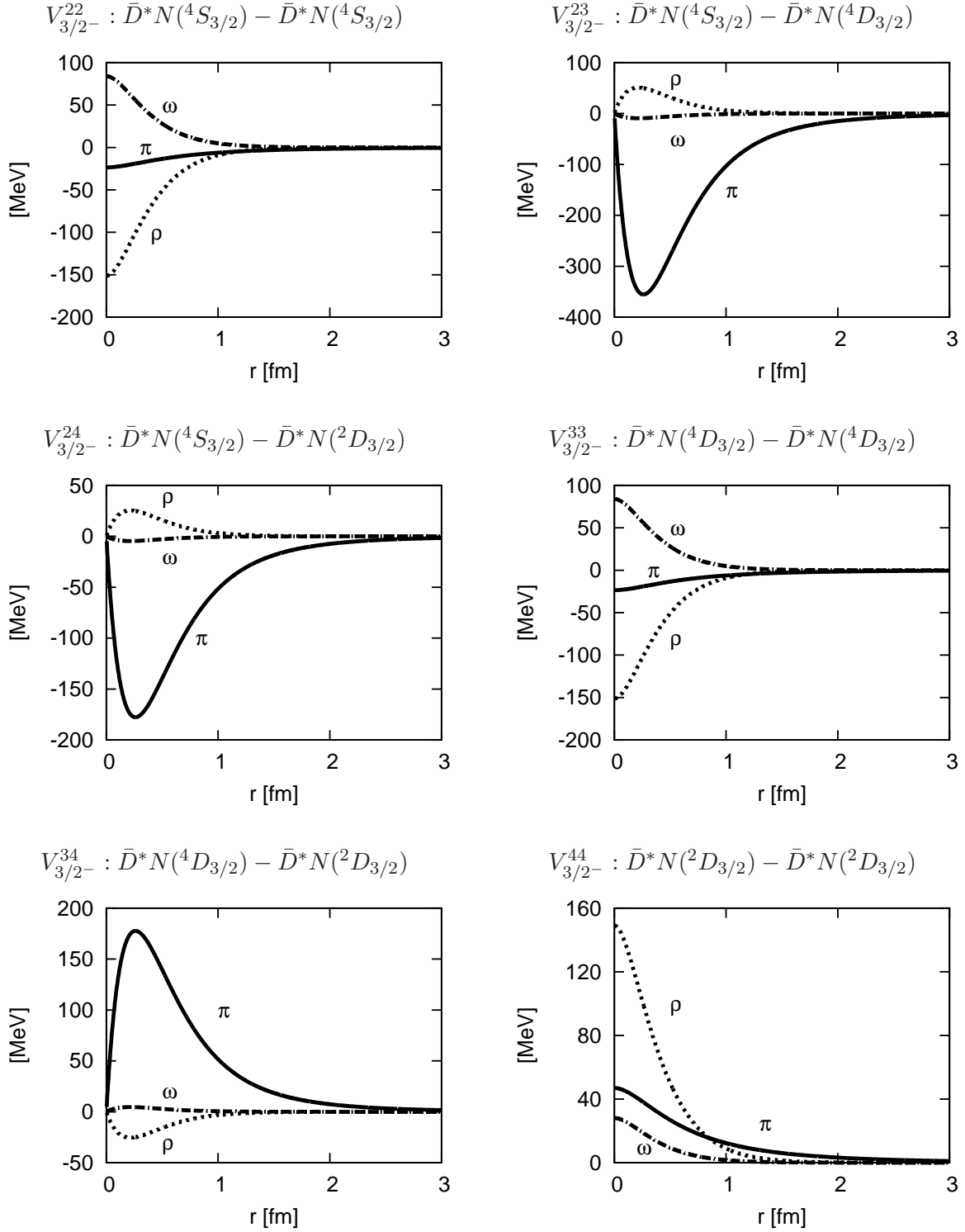


FIG. 12: Continued from Fig. 11.

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